

Dirac Gauginos: A User Manual

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Abstract

The issue of a Majorana, Dirac or pseudo-Dirac mass for gauginos must not be reduced to a question of an extension of the Minimal Supersymmetric Standard Model by extra states, parameters and phenomenological implications. On the contrary, it is intimately related to the fundamental issue of the realization of new symmetries in nature, R -symmetries. We present here a very dense compilation of the main features of models with (pseudo-)Dirac gauginos.

1 Introduction: identification of R -symmetry breaking terms

Symmetries rule the fundamental law of nature. They provide the guideline for building theoretical models designed to describe the world. While two degrees of freedom are sufficient to describe a neutral massive fermion, four degrees of freedom are necessary if the fermion carries charges under a conserved symmetry. More precisely, two elementary Majorana fermions have then to pair up to produce a Dirac fermion with the same mass.

If supersymmetry is discovered at collider experiments, there is no doubt that the question of the nature of gauginos masses will then become of central importance. This is because going from Majorana to Dirac or pseudo-Dirac masses is not just a question of extra states and parameters but also a signal the presence of R -symmetry, at least in a sector of the theory. R -symmetries are special as they do not commute with the supersymmetry supercharges, and thus acts differently on the components of superfields.

In the framework of global supersymmetric models considered here, R -symmetry appears as a continuous symmetry (which can be broken to a discrete subgroup). It can not be spontaneously broken at the electroweak scale with a generic vacuum expectation value (vev), as this would lead to a massless R -axion with a coupling insufficiently suppressed to have evaded early discovery. There remain two options: either it is conserved or explicitly broken.

In order to quantify the required size of R -symmetry breaking, one needs to identify the minimal set of operators that violate the symmetry.

1. First, it is usual to consider that R -symmetry is broken in the Minimal Supersymmetric extension of the Standard Model (MSSM) by Majorana gaugino masses. We can however use instead Dirac masses for the gauginos, pairing them with additional states in adjoint representations: a singlet \mathbf{S} for $U(1)_Y$, a triplet \mathbf{T} for $SU(2)_w$ and an octet \mathbf{O}_g for $SU(3)_c$, and preserve R -symmetry.
2. Second, in $N = 1$ supergravity, the gravitino mass required in flat space-time breaks R -symmetry. Again, this can be avoided by giving a Dirac mass for the gravitino. This requires the gravitational multiplet to be in extended $N = 2$ representations. To illustrate such a scenario, consider that the $N = 1$ gauge and matter fields appear on 3-branes. These are localized in a bulk having one flat extra dimension of radius R . Then a Dirac gravitino mass of size $1/2R$, and preserving R -symmetry, is obtained when the $N = 2$ supergravity is broken through a Scherk-Schwarz mechanism.
3. Finally, the presence of the μ and $B\mu$ mass terms for the Higgs doublets is incompatible with an R -symmetry. This makes it difficult to arrange a satisfactory electroweak symmetry breaking which preserves R -symmetry. We shall therefore consider the Higgs sector as the main source of R -symmetry breaking.

2 Dirac versus Majorana gauginos

Let us point out some of the main features and challenges raised by Dirac gauginos.

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- **Minimal content:** The adjoint representation \mathbf{S} , \mathbf{T} and \mathbf{O}_g represent the minimal set of superfields to be added to the MSSM to allow (pseudo-)Dirac masses. Such states can be remnants of an $N = 2$ extended supersymmetric structure at higher energies, that is broken by interactions with the $N = 1$ matter sector [1]. Therefore, an experimental discovery of such new states might be viewed as an indication for the extended supersymmetry origin of the model.
- **Softness:** As Majorana masses, Dirac ones can be soft [1]. However, to achieve this important property, one needs the presence of new trilinear terms [1, 2]. In fact, Dirac gaugino masses can be described with superfields by the Lagrangian:

$$\mathcal{L}_{gaugino}^{Dirac} = \int d^2\theta \left[\sqrt{2}\mathbf{m}_{1D}^\alpha \mathbf{W}_{1\alpha} \mathbf{S} + 2\sqrt{2}\mathbf{m}_{2D}^\alpha \text{tr}(\mathbf{W}_{2\alpha} \mathbf{T}) + 2\sqrt{2}\mathbf{m}_{3D}^\alpha \text{tr}(\mathbf{W}_{3\alpha} \mathbf{O}_g) \right] + h.c. \quad (2.1)$$

where we have introduced spurion superfields $\mathbf{m}_{\alpha i D} = \theta_\alpha m_{i D}$. The integration over the Grassmannian coordinates leads then to

$$\int d^2\theta 2\sqrt{2}m_D \theta^\alpha \text{tr}(W_\alpha \Sigma) \supset -m_D(\lambda_a \psi_a) + \sqrt{2}m_D \Sigma_a D_a \quad (2.2)$$

Then with $D_b^a = -g_b \phi_i^\dagger R_b^a(i) \phi_i$ (where $R_b^a(i)$ is the a^{th} generator of the group b in the representation of field i , and $R_Y^b(i) = Y(i)$ for the hypercharge) we find $\mathcal{L} \supset -m_D \sqrt{2} g_b \Sigma_a \phi^\dagger R_b^a \phi$. The new terms are proportional to the Dirac masses.

Simple scaling arguments indicate that such soft terms are in fact even softer than Majorana ones; they lead to finite contributions i.e. without logarithmic divergences [3]. Note however, that in realistic models, as the Higgs sector induces a small R -symmetry breaking, higher orders radiative corrections might lead to resurgence of logarithmic divergences.

- **Unification:** New states can be introduced at some mass scale to modify at will the running of gauge couplings, if one allows their number and charges to remain arbitrary. Though it can no more be viewed as a prediction, unification of couplings can be engineered. The non-singlet extra states have dramatic effects on the running of the standard model gauge couplings. Already, the adjoint octet spoils the asymptotic freedom of the strong interaction, making the one-loop beta function coefficient vanish and raises a difficulty. In order to keep the quantitative predictability of the theory, it is essential to ensure that the couplings remain perturbative in all the energy range where the theory is studied. This puts strong constraints on additional states that are needed to achieve unification. It makes not easy to simultaneously provide non-gravitational messengers of supersymmetry breaking. Still, such a possibility can be realized as shown in [4].
- **Generating gaugino masses:** In a top-down approach, one hopes that the gaugino masses can originate from dynamical breaking that leads to either D -terms $\langle D \rangle \neq 0$, R -preserving F -terms $\langle F \rangle \neq 0$ or both. This would allow to shed some light on the origin of hierarchies in the sparticles spectrum.

In a gauge mediation type scenario, new states are introduced at a mass scale M_{mess} to serve as mediators of the breaking, and allow to generate Dirac gaugino masses. The gaugino masses are of order [5]:

$$m_D^{(\langle D \rangle)} \sim \frac{\langle D \rangle}{M_{mess}} \quad \text{and} \quad m_D^{(\langle F \rangle)} \sim \frac{\langle F \rangle^2}{M_{mess}^3} \quad (2.3)$$

with the corresponding one-loop factors. Note that to obtain sizable contributions to all soft terms from the second case, the scale of the F -term needs to be close to the messengers scale, therefore both low. This leads to a loss of perturbative nature of the gauge couplings at intermediate scales [6, 7]. It is therefore preferable to rely on D -term contributions. In which case, the Dirac mass can be more easily extracted from computing the appropriate $U(1)$ mixing kinetic functions [8]. Note that the same messenger states can be used to restore unification of the gauge couplings [4].

In gravity mediation scheme, one needs to consider models where the spontaneous supersymmetry breaking results in a Dirac gravitino of mass $m_{3/2}$. The mediation of the breaking to the observable sector gauginos can not generate Majorana gaugino masses as they are protected by R -symmetry, but can induce Dirac masses. From dimensional analysis, one can easily see that the expected size for gaugino and scalar masses is generically to be of order $m_{3/2}$.

As a special case, let us mention the possibility of a set-up, where the Dirac gravitino is the messenger that communicates supersymmetry breaking to the observable sector through radiative effects. The Dirac gaugino mass is now generated at two-loops and found to be of order $\frac{m_{3/2}^2}{M_{Plank}}$. The power suppression (square) arises from the fact that only one of the two degenerate Majorana gravitinos interact with the observable sector. The scalar soft terms are also of the same order in $m_{3/2}$.

- **Adjoint scalars soft masses:** The generation of a soft mass for the adjoint scalars turns out to be less trivial than one expects in the case of gauge mediation. This is because the simplest interactions between the DG-adjoints and the messengers, as the Yukawa couplings descending from $N = 2$ supersymmetric gauge structure, lead to tachyonic masses. Historically, this was the main reason for abandoning the Dirac gaugino scenario in [1]. To avoid such a result, the required forms of the adjoint-messengers interactions have been fully classified in [5].

Two other important constraints need to be considered when dealing with these masses. One is that the $U(1)$ adjoint scalar is a singlet, and therefore subject to more stringent constraints, as to avoid unwanted tadpoles [4]. The second, is that the $SU(2)_w$ adjoint needs to be sufficiently massive to avoid getting a large vev after electroweak symmetry breaking.

- **Scalars soft masses:** At leading order, at which R -symmetry breaking effects are neglected, the scalar masses appear to be finite [3]. They are induced by radiative corrections from Dirac gaugino masses. This feature has been used in [9] to partially ease the supersymmetric flavor problem, as it allows to keep radiatively stable a desired hierarchy between sparticles.

Note that, contrary to the gaugino case, it is now the D -term contribution to the scalar masses that appears to be dangerously small, as it might lead to a too light selectron, for example. Therefore, all together, a preferable pattern for model building is to combine the D -term induce gaugino mass with F -term induced scalar masses [4, 10].

- **RGE running of masses:** The renormalization group equations (RGE) for running masses are different from the MSSM ones. Contrary to the Majorana case, the Dirac gaugino masses do not enter at one-loop in the RGE of squarks and sleptons [8], leading to new patterns in the generated spectra.
- **Electroweak sector:** Our simple model for the Higgs sector has a superpotential

$$W^{(s)} = \mu \mathbf{H}_u \cdot \mathbf{H}_d + \lambda_S \mathbf{S} \mathbf{H}_u \cdot \mathbf{H}_d + 2\lambda_T \mathbf{H}_d \cdot \mathbf{T} \mathbf{H}_u + \frac{M_S}{2} \mathbf{S}^2 + \frac{\kappa}{3} \mathbf{S}^3 + M_T \text{tr}(\mathbf{T} \mathbf{T}) \quad (2.4)$$

with supersymmetry breaking described by the soft masses

$$\begin{aligned} -\Delta \mathcal{L}^{soft} &= m_{H_u}^2 |H_u|^2 + m_{H_d}^2 |H_d|^2 + [\tilde{B} \mu H_u \cdot H_d + h.c.] \\ &+ m_S^2 |S|^2 + \frac{1}{2} B_S (S^2 + h.c.) + 2m_T^2 \text{tr}(T^\dagger T) + B_T (\text{tr}(TT) + h.c.) \\ &+ A_S \lambda_S S H_u \cdot H_d + 2A_T \lambda_T H_d \cdot T H_u + \frac{1}{3} \kappa A_\kappa S^3 \end{aligned} \quad (2.5)$$

where as in [11, 12] we have chosen to allow R -symmetry to be broken by a $B\mu$ mass. In addition, we have included another source of R -symmetry through $\kappa \neq 0$ term.

The singlet cubic interaction can be used to generate an effective μ -term $\tilde{\mu}$ through the singlet vev. As κ breaks the R -symmetry, it induces both a $B\mu$ term of order $\tilde{B}\mu \simeq \frac{\kappa}{\lambda_S} \tilde{\mu}^2$ as well as a singlino Majorana mass of order $M'_1 \simeq \frac{2\kappa}{\lambda_S} \tilde{\mu}$.

An important feature of the new electroweak sector is the presence of a term $\lambda_S \mathbf{SH}_u \cdot \mathbf{H}_d$ as stressed in [12]. This allows one to increase the tree-level mass of the lightest Higgs and be above the LEP bound without appealing to the large stop-top loop corrections [13].

We remind that to satisfy the constraints from by the electroweak precision tests, the triplet vev must remain small. This requires the scalar triplet mass to be an order of magnitude larger than the electroweak scale, which can be achieved easily in models of gauge mediation [4].

- **Collider signatures:** The main collider signature is the production of scalar octets at LHC. It requires the octet mass to be of the order of 1 TeV and needs use of peculiar cuts in analysis of LHC data, as described in [14].
- **Dark matter:** On one hand, the LSP candidate can be a neutralino. Unless there is a remanent $N = 2$ structure at low energies, this is a Majorana particle [15] nearly degenerate with the state made by the NLSP, which opens up new co-annihilation channels that help to get the correct dark matter relic abundance. On the other hand, the LSP can be the gravitino. When this appears as a Dirac or pseudo-Dirac state, the usual computation of the relic abundance is modified by the interactions of the “hidden Majorana gravitino” both with the observable one and with the (hidden) matter fields [15, 16].

Acknowledgement

I am grateful to the organizers of the 10th Hellenic School on Elementary Particle Physics and Gravity, Corfu 2010 for their kind hospitality and to I. Antoniadis, A. Delgado, M. D. Goodsell and M. Quir s for illuminating discussions during collaborations on different parts of the material presented in these proceedings. This work is supported in part by the European contract “UNILHC” PITN-GA-2009-237920.

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